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Weak $\Lambda N \to NN$ Transition in the Direct Quark Mechanism*

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The weak $\Lambda N \to NN$ transition is studied in the valence quark model approach. The momentum transfer for this transition is so large that the short-distance two baryon dynamics must be taken into account. The two baryon system is described in the quark cluster model and the weak transition amplitude is calculated by evaluating the matrix elements of the effective weak $\Delta S = 1$ hamiltonian. The results indicate some qualitative differences when compared with those in conventional meson-exchange calculations. Especially, we conclude that contributions of the $\Delta I = \frac{3}{2}$ transition are significant and that the discrepancy in the $n - p$ ratio between theory and experiment could be resolved by including the direct-quark processes.

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1. Introduction

The hyperon Λ decays weakly into a nucleon and a pion in the free space. Two isospin modes, $p\pi^-$ and $n\pi^0$, share 64% and 36% of the total decay. If the decay goes through a $\Delta I = \frac{1}{2}$ vertex, the $p\pi^-/n\pi^0$ ratio would be two to one except for a small correction due to the phase space difference. The experimental ratio is very close to the $\Delta I = \frac{1}{2}$ prediction, and thus support the $\Delta I = \frac{1}{2}$ hypothesis for the hadronic weak decay.

In the nuclear medium, the $\Lambda \to N\pi$ decay is suppressed by the Pauli blocking on the final nucleon state, whose momentum is less than 100 MeV/c for the Λ decay at rest. Indeed, in heavy hypernuclei, the decay is predominantly the nonmesonic one, that is, $\Lambda N \to NN$. If we assume that the initial Λ and the nucleon are at rest, then the final relative momentum of NN is about 420 MeV/c and thus is well above the Fermi momentum.

The purpose of this report is to study the direct quark processes in the two-body $\Lambda N \to NN$ weak decays and to show qualitative differences from the conventional picture employing the meson exchange mechanism. We present a possibility to solve the problems that the meson exchange mechanism encounters.

The nonmesonic decays of hypernuclei seem quite useful in studying the low energy nonleptonic weak interactions among quarks. The final relative momentum of NN is high enough to look into the short-distance component of the two nucleon system. This type of the hyperon decay may reveal a new aspect of the weak interaction under the influence of the strong interaction. An advantage of using the hypernuclear decay is that the process is selective in isospin, spin and orbital angular momentum for appropriate initial and final states of the hypernucleus.

Theoretical study of the nonmesonic decay of hypernuclei has traditionally employed the meson (π, K, ρ, etc.) exchange mechanism, where one of the meson-baryon vertices involves the weak transition $s \to d$[1]. Accumulating experimental data, however, have revealed some difficulties in the meson-exchange picture. For instance, the so-called $n-p$ ratio, i.e., the ratio R_{np} of $\Lambda n \to nn$ v.s. $\Lambda p \to np$ decay in the nucleus, is predicted very small, $R_{np} \simeq 0.1 - 0.4$ in the meson-exchange picture. This is due to the strong

contribution of the tensor force, which is preferred at the large momentum transfer. The tensor force selects the $S = 1$, $I = 0$ pn final state and therefore R_{np} becomes small. The experimental data seem not to agree with the prediction, i.e., $R_{np}^{exp} \simeq 1$ in decays of light hypernuclei. We argue that the direct quark process, which does not follow the $I = 0$ selection rule, may enhance the $n-p$ ratio.

The mesonic weak decays of hyperons have been tested for the $\Delta I = \frac{1}{2}$ rule and are known to satisfy the rule to about 5% error. The same rule for the nonmesonic weak processes, like $\Lambda N \to NN$, is not confirmed yet. Indeed, an analysis of the decay of the $A = 3$ and 4 hypernuclei claims that the $\Delta I = \frac{1}{2}$ rule may be violated[2]. It is therefore urgent to clarify the mechanism of the $\Delta I = \frac{1}{2}$ rule in the free hyperon decays and to study whether the same mechanism restricts the nonmesonic decays to $\Delta I = \frac{1}{2}$ as well.

In the study of the meson-exchange processes, the $\Delta I = \frac{1}{2}$ rule is assumed from the beginning, implemented in the $\Lambda \to N\pi$ vertex. We instead employ the effective quark-quark weak hamiltonian, which contains both the $\Delta I = \frac{1}{2}$ and $\Delta I = \frac{3}{2}$ components. Although the $\Delta I = \frac{3}{2}$ part has a small overall coefficient, we will see that its matrix elements for the $\Lambda N \to NN$ decay may not be small compared to the $\Delta I = \frac{1}{2}$ component.

The paper is organized as follows. Sect 2 is devoted to the basic formulation of the present calculation. The effective hamiltonian and the quark cluster model wave functions are presented. Various approximations employed in this calculation are examined. We present the results of the calculation in sect. 3. They are used for studying qualitative differences between the direct-quark mechanism and the conventional meson-exchange picture. A conclusion is given in sect. 4.

2. Formulation

2.1. Effective Weak Hamiltonian

The effective weak hamiltonian describing $\Delta S = 1$ processes has been calculated by several authors[3,4,5]. It can be computed by analyzing the correction due to the strong

interaction on the pure weak vertex $su \to du$:

$$H(\text{purely weak:}\Delta S=1) = -\frac{G_f}{\sqrt{2}} \sin\theta_c\, (\bar{u}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta u_\beta)_{V-A} \tag{1}$$

where

$$(\bar{u}_\alpha s_\alpha)_{V-A} \equiv (\bar{u}_\alpha \gamma^\mu (1-\gamma_5)\, s_\alpha) \quad etc. \tag{2}$$

and α and β denote the color of quarks and the color sum is always assumed. Note that there exists no strangeness-changing neutral current in the standard electro-weak theory, and thus the vertex is only for the left-handed quarks.

This purely weak four-quark vertex clearly contains both the $\Delta I = \frac{1}{2}$ and $\Delta I = \frac{3}{2}$ components. It was pointed out, however, that the correction due to the strong interaction enhances the $\Delta I = \frac{1}{2}$ component while the $\Delta I = \frac{3}{2}$ is suppressed at the same time[3]. The strong correction is treated perturbatively at the scale $Q^2 \simeq M_W^2$, and only the lowest order diagrams are taken into account. The mechanism of the $\Delta I = \frac{1}{2}$ enhancement can be understood by realizing the anomalous dimensions of the $\Delta I = \frac{1}{2}$ and $\Delta I = \frac{3}{2}$ components in eq.(1) have opposite signs with each other. The $\Delta I = \frac{1}{2}$ ($\Delta I = \frac{3}{2}$) anomalous dimension is positive (negative) and therefore, when the renormalization scale is moved down from the W boson mass to the low-energy hadronic scale ($\lesssim$1 GeV), the operator with the positive anomalous dimension is enhanced and vice versa. The renormalization group equation also induces new four-quark operators through operator mixings. Another contribution comes from the so-called penguin diagrams. They are purely $\Delta I = \frac{1}{2}$ and thus help the $\Delta I = \frac{1}{2}$ rule.

Taking these effects into account, the low energy effective weak hamiltonian has been derived[5]:

$$H_{eff}^{\Delta S=1}\left(Q^2 \sim \mu^2\right) = -\frac{G_f}{\sqrt{2}} \sum_{r=1, r\neq 4}^{6} K_r O_r \tag{3}$$

where the four-quark operators, O_k ($k=1$, 2, 3, 5 and 6) are defined by[3]

$$O_1 = (\bar{d}_\alpha s_\alpha)_{V-A} (\bar{u}_\beta u_\beta)_{V-A} - (\bar{u}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta u_\beta)_{V-A} \tag{4}$$

$$O_2 = (\bar{d}_\alpha s_\alpha)_{V-A} (\bar{u}_\beta u_\beta)_{V-A} + (\bar{u}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta u_\beta)_{V-A}$$

$$
\begin{aligned}
&+\ 2(\bar{d}_\alpha s_\alpha)_{V-A}(\bar{d}_\beta d_\beta)_{V-A} + 2(\bar{d}_\alpha s_\alpha)_{V-A}(\bar{s}_\beta s_\beta)_{V-A} && (5)\\
O_3 &= O_3(\Delta I = \frac{1}{2}) + O_3(\Delta I = \frac{3}{2}) && (6)\\
O_3(\ \Delta\ I = \frac{1}{2}) &= \frac{1}{3}\Big[(\bar{d}_\alpha s_\alpha)_{V-A}(\bar{u}_\beta u_\beta)_{V-A} + (\bar{u}_\alpha s_\alpha)_{V-A}(\bar{d}_\beta u_\beta)_{V-A}\\
&+\ 2(\bar{d}_\alpha s_\alpha)_{V-A}(\bar{d}_\beta d_\beta)_{V-A} - 3(\bar{d}_\alpha s_\alpha)_{V-A}(\bar{s}_\beta s_\beta)_{V-A}\Big]\\
O_3(\ \Delta\ I = \frac{3}{2}) &= \frac{5}{3}\Big[(\bar{d}_\alpha s_\alpha)_{V-A}(\bar{u}_\beta u_\beta)_{V-A} + (\bar{u}_\alpha s_\alpha)_{V-A}(\bar{d}_\beta u_\beta)_{V-A}\\
&-\ (\bar{d}_\alpha s_\alpha)_{V-A}(\bar{d}_\beta d_\beta)_{V-A}\Big]\\
O_5 &= (\bar{d}_\alpha s_\alpha)_{V-A}(\bar{u}_\beta u_\beta + \bar{d}_\beta d_\beta + \bar{s}_\beta s_\beta)_{V+A} && (7)\\
O_6 &= (\bar{d}_\alpha s_\beta)_{V-A}(\bar{u}_\beta u_\alpha + \bar{d}_\beta d_\alpha + \bar{s}_\beta s_\alpha)_{V+A} && (8)
\end{aligned}
$$

Among these operators, O_3 contains a part that induces the $\Delta I = \frac{3}{2}$ transition, $O_3(\Delta I = \frac{3}{2})$, while the others are purely $\Delta I = \frac{1}{2}$. The coefficients K_r can be calculated by solving the renormalization group equation to the one-loop QCD corrections. They depend on the mass of the top quark m_t through the penguin diagrams, which generate the operators O_5 and O_6 with the $V + A$ coupling. We find that the final results are insensitive to the choice of m_t and here choose $m_t = 200$ GeV/c^2. The coefficients also depend on the energy scale μ^2 for the effective hamiltonian. In the present calculation, we choose two sets of the values given in ref.[5]: $\mu = 0.24$ GeV and 0.71 GeV. They are chosen so as to give $\alpha_s(\mu^2) = 1$ for the QCD Λ_{QCD} parameter, $\Lambda_{QCD} = 0.1$ GeV and 0.316 GeV, respectively. The values of the coefficients K_r used in the present calculation are given in Table 1. One sees that the two choices are not much different except for K_5 and K_6. We will find that the differences in the transition matrix elements for these choices are at most 10%.

The most prominent feature of this effective hamiltonian is that the QCD correction enhances the O_1 component while the other terms are suppressed. This is the main mechanism for the $\Delta I = \frac{1}{2}$ enhancement as is explained above. Later we will compare the results with and without $O_3(\Delta I = \frac{3}{2})$ in order to study the $\Delta I = \frac{3}{2}$ contribution.

This effective hamiltonian has been used for the calculations of the nonleptonic decay of strange mesons and baryons[4]. It is found that although the $\Delta I = \frac{1}{2}$ transition is indeed enhanced in those decays, the enhancement is not enough to account for the experimental data quantitatively. It was suggested[7] that an additional $\Delta I = \frac{1}{2}$ enhancement arises

Table 1: Two choices of strengths of the weak effective four-fermi vertices, taken from ref.[5]. We use the version with flavor-number dependent Λ and $m_t = 200$GeV. The values of the CKM matrix elements are taken as the central values of those given in ref.[6].

	μ (GeV)	$\Lambda^{(4)}$ (GeV)	K_1	K_2	K_3	K_5	K_6
I	0.24	0.10	-0.284	0.009	0.026	0.004	-0.021
II	0.71	0.316	-0.270	0.011	0.027	0.002	-0.010

from the mesonic correction in the chiral effective theory. It was also suggested that the decay amplitudes may be sensitive to the meson and baryon wave functions[8].

2.2. Six-Quark Wave Function

In calculating the decay amplitude for $\Lambda N \to NN$, we employ the constituent quark model, which describes the spin-flavor structure of the ground-state baryons very well. Two-baryon systems are expressed by the quark-cluster-model wave functions[9,10]. First, we assume that the baryon consists of three valence quarks, whose orbital wave function is a harmonic oscillator eigenstate,

$$\phi(1,2,3)^{\text{orb}} = \left(\frac{1}{2\pi b^2}\right)^{\frac{3}{4}} \left(\frac{2}{3\pi b^2}\right)^{\frac{3}{4}} \exp\left\{-\frac{1}{4b^2}\vec{\xi}_{12}^{\,2}\right\} \exp\left\{-\frac{1}{3b^2}\vec{\xi}_{12-3}^{\,2}\right\} \tag{9}$$

where ξ's are the Jacobi coodinates and the Gaussian parameter b is chosen $b = 0.5$ fm. We here neglect the asymmetry due to the mass difference of the s and u quarks in the Λ wave function. The six quark wave functions are given by

$$\begin{aligned} |\Lambda N\rangle &= \mathcal{A}^6 |\phi(1,2,3)\phi(4,5,6)\chi_0(\vec{R})\rangle \\ |NN\rangle &= \mathcal{A}^6 |\phi(1,2,3)\phi(4,5,6)\chi(\vec{R})\rangle \end{aligned} \tag{10}$$

where $\mathcal{A}^6$ is the antisymmetrization operator for six quarks, ϕ is the internal wave function of the baryon, and $\vec{R}$ is the relative coodinate of two baryons. $\chi_0(\vec{R})$ $(\chi(\vec{R}))$ is the initial (final) relative wave function. The flavor-spin part of ϕ is taken to be purely the SU(6) wave function.

In the present calculation, we choose the simplest relative wave functions, *i.e.*, a Gaussian for the initial state and the plane wave for the final state.

$$\chi_0(\vec{R}) = \left(\frac{1}{\pi B^2}\right)^{\frac{3}{4}} \exp\left\{-\frac{1}{2B^2}\vec{R}^2\right\} \tag{11}$$

$$\chi(\vec{R}) = \exp\left\{i\vec{k}\cdot\vec{R}\right\} \tag{12}$$

The relative momentum of the final state is determined by the realistic Q value $\Delta E \equiv M_\Lambda - M_N$ of the decay: $k = 416$ MeV/c. The Gaussian B parameter of the initial state is to be determined by the ΛN wave function in the hypernucleus. Suppose that we choose (unrealistically) $B = \sqrt{2/3}\,b$, then the initial wave function is reduced to the $(0s)^6$ configuration in the harmonic oscillator shell model. This can be interpreted as a dibaryon state in which the initial Λ and N are on top of each other ($R = 0$). Later we show the results of the calculation where we employ $B = 1.838$ fm ($= \sqrt{2}\times$ 1.3), that corresponds to a ΛN system in the ^{4}He nucleus.

Although these choices of the wave functions may not be totally realistic, they will clarify the qualitative difference between the meson-exchange and the direct-quark processes, which is the purpose of this study. An advanced calculation using the realistic two-baryon wave functions is under way.

Because we employ the nonrelativistic valence quark picture for the wave functions, the effective hamiltonian is also approximated by adopting the Breit-Fermi nonrelativistic expansion up to p/m. Then the spin-orbital part of the operator contains terms which conserve parity,

$$1\times 1, \quad (\vec{\sigma}_i\cdot\vec{\sigma}_j)\times 1$$

and those which break parity,

$$(\vec{\sigma}_i\pm\vec{\sigma}_j)\cdot\vec{q}_{ij}, \quad (\vec{\sigma}_i\pm\vec{\sigma}_j)\cdot\left(\vec{P}_i\pm\vec{P}_j\right), \quad i\,(\vec{\sigma}_i\times\vec{\sigma}_j)\cdot\vec{q}_{ij}, \quad i\,(\vec{\sigma}_i\times\vec{\sigma}_j)\cdot\left(\vec{P}_i\pm\vec{P}_j\right)$$

where $\vec{q}_{ij} = \vec{p}_i^{\,\prime} - \vec{p}_i$ and $\vec{P}_i = \frac{1}{2}(\vec{p}_i + \vec{p}_i^{\,\prime})$ with $\vec{p}_i$ ($\vec{p}_i^{\,\prime}$) the initial (final) momentum of the i-th quark. In this expansion, we take account of the SU(3) breaking effects due to the quark mass differences. Then the coefficients for these operators depend on the light quark constituent mass, $m_q = m_u = 313$ MeV and the mass ratio, $m_u/m_s = 0.6$.

Table 2: Possible initial and final quantum numbers for the initial $L = 0$ transition

channel	isospin	spin–orbital	
1	$p\Lambda \to pn$	${}^1S_0 \to {}^1S_0$	a_p
2		${}^1S_0 \to {}^3P_0$	b_p
3		${}^3S_1 \to {}^3S_1$	c_p
4		${}^3S_1 \to {}^3D_1$	d_p
5		${}^3S_1 \to {}^1P_1$	e_p
6		${}^3S_1 \to {}^3P_1$	f_p
7	$n\Lambda \to nn$	${}^1S_0 \to {}^1S_0$	a_n
8		${}^1S_0 \to {}^3P_0$	b_n
9		${}^3S_1 \to {}^3P_1$	f_n

In the present study, we restrict our initial state to $L = 0$. Table 2 shows nine possible combinations of L, S, J, and I for the initial and final states. We note that the $I = 1$ final states are allowed both for $(\Lambda n \to nn)$ and $(\Lambda p \to pn)$, while the $I = 0$ states are not possible for $(\Lambda n \to nn)$. Thus we have 6 $(\Lambda p \to pn)$ and 3 $(\Lambda n \to nn)$ matrix elements, which are labeled from a through f in Table 2, according to the widely used notation[11]. Among them, the channels 1, 3, 4, and 7 (a, c and d) are the parity conserving transitions, while the others violate the parity invariance. The amplitude d is not zero only for the tensor component of the weak interaction, which we neglect in our quark model calculation by truncating the p/m expansion.

The transition amplitudes are calculated according to the standard quark cluster model approach. Remember that we take into account the full antisymmetrization among six quarks. One needs to calculate the exchange matrix elements as well as the direct ones.

3. Results

The results of the calculation are summarized in Table 3. Nine amplitudes give all the information for the $\Lambda N \to NN$ weak decay from $L = 0$. We compare the results for

Table 3: Calculated transition matrix elements in 10^{-10} $MeV^{-1/2}$. The numbers without parenthesis (in parenthesis) are the results for the parameter I (II).

channel	full		$\Delta I = \frac{3}{2}$ omitted		OPE
a_p	-6.66	(-6.83)	-0.02	(0.04)	4.52
b_p	5.79	(6.30)	0.06	(0.37)	-24.4
c_p	2.70	(2.48)	2.70	(2.48)	4.52
d_p	0	(0)	0	(0)	-83.0
e_p	-2.21	(-2.01)	-2.21	(-2.01)	-42.3
f_p	-5.57	(-5.02)	-5.39	(-4.83)	19.9
a_n	4.67	(4.92)	-0.03	(0.06)	6.39
b_n	-3.96	(-3.67)	0.09	(0.52)	-34.5
f_n	-7.49	(-6.70)	-7.62	(-6.83)	28.2

the two choices of the weak hamiltonian parameters, which differ in the energy scale or Λ_{QCD}, given in Table 1 and see that their differences are small. The numbers given under the "$\Delta I = \frac{3}{2}$ omitted" are the results without the $\Delta I = \frac{3}{2}$ component of the O_3 operator. In this case, the ratio a_n/a_p, b_n/b_p, and f_n/f_p are equal to $\sqrt{2}$. The other amplitudes, c, d and e do not contain any $\Delta I = \frac{3}{2}$ component, because the final NN states have $I = 0$. We find that the amplitudes a and b get significant contributions from the $\Delta I = \frac{3}{2}$ component, while f has only a small contribution of $\Delta I = \frac{3}{2}$. Thus we conclude that the $\Delta I = \frac{3}{2}$ transition can be studied in the weak decay starting from the ΛN ${}^1S_0$ state.

In Table 4, we summarize the calculated decay rates with the initial spin averaged and the final states summed up.

$$\Gamma_p = \frac{\pi M_N k}{(2\pi)^3} \frac{1}{4} [a_p^2 + b_p^2 + 3(c_p^2 + d_p^2 + e_p^2 + f_p^2)] \tag{13}$$

$$\Gamma_n = \frac{\pi M_N k}{(2\pi)^3} \frac{1}{4} [a_n^2 + b_n^2 + 3f_n^2] \tag{14}$$

We again find that the $\Delta I = \frac{3}{2}$ component of O_3 changes the total decay rate by as much as a factor two for the direct quark processes.

Table 4: Calculated observables. The numbers without parenthesis (in parenthesis) are the results for the parameter I (II).

	full		$\Delta I = \frac{3}{2}$ omitted		OPE
Γ_p (10^7 sec^{-1})	0.39	(0.36)	0.23	(0.19)	52.4
Γ_n (10^7 sec^{-1})	0.39	(0.32)	0.33	(0.26)	6.8
R_{np}	0.99	(0.89)	1.41	(1.39)	0.13
η_p	2.13	(1.96)	4.64	(4.47)	0.34
η_n	8.43	(6.12)	$2.23{\cdot}10^5$	$(4.04{\cdot}10^4)$	87.5
$a_1(p)$	−0.36	(−0.32)	−0.58	(−0.58)	−0.19

The $n-p$ ratio,

$$R_{np} \equiv \Gamma_n/\Gamma_p, \tag{15}$$

the ratio of the parity violating (PV) v.s. the parity conserving (PC) contributions,

$$\eta_p = \frac{b_p^2 + 3(f_p^2 + e_p^2)}{a_p^2 + 3(c_p^2 + d_p^2)} \tag{16}$$

$$\eta_n = \frac{b_n^2 + 3f_n^2}{a_n^2} \tag{17}$$

and the decay asymmetry parameter, a_1, are also given in Table 4. We find that the $n-p$ ratio for the direct quark process is much larger than that obtained in the meson-exchange calculation. The pure $\Delta I = \frac{1}{2}$ calculation also yields large R_{np}, indicating that the R_{np} enhancement is not related to the $\Delta I = \frac{1}{2}$ rule violation.

The results of the one-pion exchange transition are also shown in Tables 3 and 4, for comparison[1]. These amplitudes satisfy $a_n/a_p = b_n/b_p = f_n/f_p = \sqrt{2}$, because $\Delta I = \frac{1}{2}$ is assumed for the weak pion-baryon vertex. One sees that the amplitude d in this case is dominant. This comes from the tensor part of the one pion exchange and is enhanced due to a large relative momentum in the final state. Because the amplitude d is not allowed for the $n\Lambda \to nn$ by the Pauli principle, the $n-p$ ratio in the pion exchange amplitudes

becomes very small. Remember that the tensor part of the direct quark interaction is neglected because it is of order $(p/m)^2$.

The magnitudes of the transition amplitudes are in general larger for the meson exchange mechanism than the direct-quark process. It is noticed, however, that some amplitudes, such as a and c, have comparable direct-quark amplitudes. Therefore, if one can select the initial and/or final spin states in the decay experiments, it will be possible to detect the contribution of the direct quark processes.

The decay asymmetry parameter describes the angular distribution of the outgoing two nucleons in the rest frame,

$$W(\theta) = 1 + a_1 \mathcal{P}_\Lambda P_1(\cos\theta) \tag{18}$$

where $\mathcal{P}_\Lambda$ is the polarization of Λ in the nucleus. The parameter a_1 is given in terms of the two-body decay amplitudes by

$$a_1 = \frac{2\sqrt{3}\,(\sqrt{2}c + d)f}{a^2 + b^2 + 3(c^2 + d^2 + e^2 + f^2)} \tag{19}$$

and thus indicates the interference between the PV and PC components. For the $\Lambda n \to nn$ decay, c and d vanish due to the isospin conservation and therefore $a_1(n)$ is zero. Recent experiment done at KEK indicates a large negative $a_1(p)$ for light hypernuclei[12]. The data is consistent with $a_1(p) \leq -0.6$. Our calculation yields the correct sign, but the magnitude is smaller. The meson exchange calculation done by Ramos *et al.*, also predicted a small magnitude with the correct sign[1].

Recently, Schumacher[2] suggested that analyses of the non-mesonic decays of the A=4 and 5 hypernuclei may be usuful in checking the $\Delta I = \frac{1}{2}$ rule for the nonmesonic weak decay. One can parametrize the n-p ratios of the decay of ${}^4_\Lambda$He and ${}^5_\Lambda$He and the ratio of non-mesonic decay widths of ${}^4_\Lambda$He and ${}^4_\Lambda$H in terms of R_{nS} and R_{pS}, where R_{NS} stands for the decay rate of ΛN with the initial spin S[11]. Then, using experimental data, the ratios of R_{NS} can be extracted. In fact, the R_{n0}/R_{p0} ratio will be 2 for the pure $\Delta I = \frac{1}{2}$ transition, while it becomes 1/2 for the pure $\Delta I = 3/2$ transition. Because these hypernuclei involve only the relative s-wave states, the direct quark processes can play significant roles. It may then be possible to see the $\Delta I = 3/2$ decay, if any. Table 5 shows

Table 5: Observables for light hypernuclei. The numbers without parenthesis (in parenthesis) are the results for the parameter I (II).

	full		$\Delta I = \frac{3}{2}$ omitted		OPE
$R_{np}({}^4_\Lambda\mathrm{He})$	0.36	(0.39)	0	(0)	0.09
$R_{np}({}^5_\Lambda\mathrm{He})$	0.99	(0.89)	1.41	(1.40)	0.13
$\Gamma_{n.m.}({}^4_\Lambda\mathrm{He})/\Gamma_{n.m.}({}^4_\Lambda\mathrm{H})$	0.78	(0.78)	0.71	(0.72)	6.26
R_{n0}/R_{p0}	0.48	(0.44)	2	(2)	2

the $J = 0$ ($S = 0$) part of the n-p ratio, R_{n0}/R_{p0} as well as the above observables for our direct quark amplitudes with and without the $\Delta I = \frac{3}{2}$ components. Our full amplitudes indeed predict a small ratio R_{n0}/R_{p0}, which indicated a large $\Delta I = 3/2$ contribution. It is therefore concluded that the $\Delta I = \frac{3}{2}$ components of the direct quark processes is possibly observed in non-mesonic decays of light hypernuclei.

4. Conclusion and Discussion

We present a quark model calculation of the direct-quark processes of the weak $\Lambda N \to NN$ decay, which can be observed exclusively in decays of hypernuclei. We find that the transition amplitudes in some of the decay channels are comparable to those in the meson exchange decays in magnitudes and show qualitatively distinctive properties. This is encouraging because the direct-quark processes may resolve the discrepancies between experiment and theory based on the meson-exchange mechanism. Indeed, after averaging over the initial and final spin states, we obtain a large $n-p$ ratio from the direct-quark amplitudes. In the actual hypernuclear decays, one expects a variety of the spin-isospin combinations. We also have to note that the initial orbital angular momentum greater than zero may contribute significantly. We, therefore, cannot make a definite prediction here. Further study of the $\Lambda N \to NN$ decay with higher partial waves and realistic two-baryon wave functions is under way.

It is also important to combine the meson-exchange amplitudes with the direct-quark ones to make a final quantitative conclusion. We here do not superpose the direct-quark and the one-pion-exchange amplitudes because both of them are yet incomplete. The one-pion exchange is far from realistic because other mesons, such as K and ρ, are known to contribute to this process significantly. It is not clear whether the phenomenological lagrangian for the mesonic decay of Λ is consistent with the weak effective hamiltonian used in the direct-quark calculation. It is important to understand the mesonic decay and its $\Delta I = \frac{1}{2}$ rule starting from the effective four-quark Lagrangian[4,7]. A study along this line is being carried out.

Our plan also includes the study of $\Lambda N \to NN$ with more realistic wave functions for the initial ΛN and the final NN states. Especially, effects of the baryon-baryon short-range correlation may change the results quantitatively. A more realistic calculation is under way.

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